

Two-Photon Transport in a Waveguide Coupled to a Cavity with a Two-level System

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We consider a system where a waveguide is coupled to a cavity embedded with a two-level system (TLS), and study the effects when a two-photon quantum state is injected into the waveguide. The wave function of two outgoing photons is exactly solved using the Lehmann-Symanzik-Zimmermann (LSZ) reduction formalism. Our results explicitly exhibit the photon blockade effects in the strong atom-cavity coupling regime. The quantum statistical characters of the outgoing photons, including the photon bunching and anti-bunching behaviors, are also investigated in both the strong and weak coupling regimes. These results agree with the observations of recent experiments.

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Introduction. — The Jaynes-Cummings (JC) system, involving cavity quantum electrodynamics (QED) for a two-level atom inside a cavity, is of most importance for quantum optics [1] and its applications. In the past decade, the JC system, in the regime where there are only a few photons inside the cavity, has been very extensively studied due to its potential applications for both quantum information processing and quantum device physics [2]. The later is usually based on some solid state systems resembling the JC system, such as the superconducting circuit QED systems [3] and optomechanical architectures [4]. Recent experiments about various JC systems include the observation of photon blockade in a macroscopic cavity coupled to an atom [5], the demonstrations of on-chip cavities coupling to atom or atom-like objects [6], as well as the integration of such atom-cavity systems with waveguides [7]. A few configurations, where the atom-cavity system is either side-coupled [8, 9], or directly coupled [5, 10] to a waveguide, are schematically shown in Fig. 1a and Fig. 1b.

In the context of these recent experiments, it is necessary to theoretically study the intensity and coherence properties of transmitted or reflected light, when a few-photon quantum state is injected into the system through a waveguide. In Ref. [5, 10], the quantum states inside the cavity were expanded on a basis of photon number states. Truncating the number of basis states then reduced the Master equations to an ordinary differential equations which can be solved numerically. This system has also been simulated by using the quantum trajectory approach [11]. Analytically, closed-form formulas regarding the transmission and coherence properties have been obtained, either in the weak excitation limit where the atom is assumed to be mostly in the ground state [12], or in a mean-field-like approach where the expectation value of operator product is taken as the product of operator expectation values [13].

In this Letter, using field-theoretic techniques [14], we provide an exact analytic formula for the out-going photon wave-functions, when a two-photon state is injected

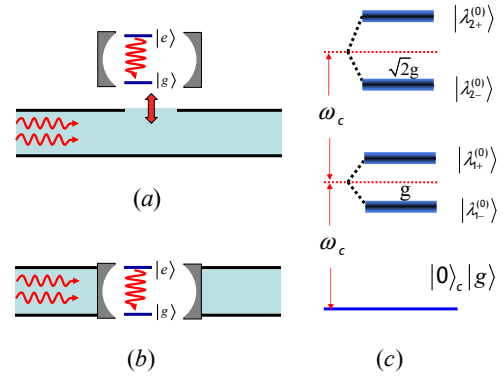


FIG. 1: (Color online) Two kinds of coupling structures and the schematic of energy spectrum. (a) The side-coupled waveguide; (b) The direct-coupled waveguide; (c) the schematic for the energy spectrum of JC model in the subspaces $n = 0, 1$, and 2 .

into the system shown in Fig. 1. Our work is in contrast to most existing theoretical works that used coherent state input. Given that these systems will be ultimately used to process quantum information, understanding their response to states other than coherent states is very important. Also, our result is exact, without the need of either the mean-field approximation, or being restricted to the weak excitation limit. In this system, the indirect photon-photon interaction is strong when the atom has significant probability in the excited state [13]. Moreover, it is known even a single-photon pulse, when properly designed, can completely invert an atom in this system [15]. Thus, it is important to go beyond the weak excitation limit.

Transport model and S -matrix. — The systems in Fig. 1a, with a waveguide side-coupled to a cavity, is described by a Hamiltonian [16] $H = H_W + H_{JC} + H_I$ containing three parts: (a) The waveguide Hamiltonian $H_W = \sum_k \varepsilon_k a_k^\dagger a_k$, where a_k (a_k^\dagger) denotes the annihilation (creation) operators of the photon. $\varepsilon_k = v|k|$ is the wave-

uide dispersion relation, and we take the speed of light v as unity. (b) The JC Hamiltonian for the coupling of the cavity field to the TLS

$$H_{JC} = \omega_c a^\dagger a + \Omega |e\rangle \langle e| + g(a^\dagger |g\rangle \langle e| + a |e\rangle \langle g|), \quad (1)$$

where a (a^\dagger) denotes the annihilation (creation) operators of the photon in the cavity with frequency ω_c , $|e\rangle$ ($|g\rangle$) denotes the excited (ground) state of the TLS with energy level spacing Ω , and g is the coupling constant of the TLS and the cavity field. (c) The term $H_I = V \sum_k (a_k^\dagger a + \text{H.c.})/\sqrt{L}$, describing the coupling between the cavity and the waveguide. This term results in the decay of the cavity mode into the waveguide. Here, V is the waveguide-cavity coupling constant and L is the length of waveguide.

In term of the bonding and anti-bonding waveguide mode operators defined as $e_k(o_k) = (a_k \pm a_{-k})/\sqrt{2}$, we can write $H = H_e + H_o$. The bonding sub-system is described by

$$H_e = \sum_{k>0} k e_k^\dagger e_k + \frac{\tilde{V}}{\sqrt{L}} \sum_{k>0} (e_k^\dagger a + \text{H.c.}) + H_{JC}, \quad (2)$$

where $\tilde{V} = \sqrt{2}V$. The anti-bonding sub-system is described by $H_o = \sum_{k>0} k o_k^\dagger o_k$ and is decoupled from the cavity.

Next, we utilize the Lehmann-Symanzik-Zimmermann reduction approach [14] to study the bonding sub-system as described by the Hamiltonian H_e . We aim to calculate the S -matrix elements between the incoming and outgoing n -photon states, specified by their photon momenta $\mathbf{k} = k_1, \dots, k_n$ and $\mathbf{p} = p_1, \dots, p_n$, respectively. The two-photon S -matrix has the form

$$S_{p_1 p_2 k_1 k_2} = S_{p_1 k_1} S_{p_2 k_2} + S_{p_2 k_1} S_{p_1 k_2} + iT_{p_1 p_2 k_1 k_2}, \quad (3)$$

where $S_{pk} = \delta_{pk} + iT_{pk}$ is the single-photon S -matrix element. To determine the T -matrix, we first exactly calculate the connected Green function $G_{[\mathbf{p};\mathbf{k}]}(\omega_{\mathbf{p}}, \omega_{\mathbf{k}}) = \int G_{[\mathbf{p};\mathbf{k}]}(\mathbf{t}', \mathbf{t}) \prod_{j=1}^n [-\exp(i\omega_{p_j} t'_j - i\omega_{k_j} t_j) dt'_j dt_j / 2\pi]$ in the frequency domain with $\omega_{\mathbf{k}} = \omega_{k_1}, \dots, \omega_{k_n}$ and $\omega_{\mathbf{p}} = \omega_{p_1}, \dots, \omega_{p_n}$, where

$$G_{[\mathbf{p};\mathbf{k}]}(\mathbf{t}', \mathbf{t}) = \frac{\delta^{2n} \ln Z[\eta_k, \eta_k^*]}{\delta \eta_{p_1}^*(t'_1) \dots \delta \eta_{p_n}^*(t'_n) \delta \eta_{k_1}(t_1) \dots \delta \eta_{k_n}(t_n)} \Big|_{\substack{\eta_k=0 \\ \eta_k^*=0}}, \quad (4)$$

with $\mathbf{t} = t_1, \dots, t_n$, and $\mathbf{t}' = t'_1, \dots, t'_n$. $Z[\eta_k, \eta_k^*] = \int D[e, a, \sigma] \exp\{i \int dt [L_e + \sum_{k>0} (\eta_k^* e_k + h.c.)]\}$ is a generating functional with L_e being the Lagrangian for the bonding subsystem, and σ denoting the variables of the TLS. The T -matrix elements can then be related to the Green function by

$$iT_{[\mathbf{p};\mathbf{k}]} = \frac{(2\pi)^n G_{[\mathbf{p};\mathbf{k}]}(\omega_{\mathbf{p}}, \omega_{\mathbf{k}})}{\prod_{j=1}^n [G_0(\omega_{p_j}, p_j) G_0(\omega_{k_j}, k_j)]} \Big|_{\substack{\omega_{p_j}=p_j \\ \omega_{k_j}=k_j}},$$

where the bare Green function G_0 defined by $iG_0^{-1}(\omega_p, p) = \omega_p - \varepsilon_p + i0^+$.

For one or two-photons, this computation results in

$$iT_{pk} = -\tilde{V}^2 \int \frac{dt' dt}{2\pi} e^{ipt' - ikt} \langle \mathcal{T} a(t') a^\dagger(t) \rangle, \quad (5)$$

and

$$iT_{p_1 p_2 k_1 k_2} = \tilde{V}^4 \int \left(\prod_{j=1,2} \frac{dt_j dt'_j}{2\pi} e^{ip_j t'_j - ik_j t_j} \right) G_4, \quad (6)$$

Here, $G_4 = \langle \mathcal{T} a(t'_1) a(t'_2) a^\dagger(t_1) a^\dagger(t_2) \rangle$ is the four point Green function of $a(t) = \exp(iH_{\text{eff}} t) a \exp(-iH_{\text{eff}} t)$, and the effective non-Hermitian Hamiltonian H_{eff} is obtained by simply replacing ω_c with $\alpha = \omega_c - i\tilde{V}^2/2$ in H_{JC} . The imaginary part of α is the cavity decay rate. $\langle \mathcal{T} \dots \rangle$ is the time-ordered average on the state $|0\rangle_c |g\rangle$. Here, $|0\rangle_c$ and $|g\rangle$ denote the vacuum state of cavity field and the ground state of TLS, respectively. Since excitation number $N = a^\dagger a + |e\rangle \langle e|$ commutes with H_{eff} , in its invariant subspace with $N = n$, H_{eff} is diagonalized with the eigenstates $|\lambda_{n\pm}\rangle = \mathcal{N}_{n\pm} \{ -\sqrt{n} g |n-1\rangle_e |e\rangle + [\Omega + (n-1)\alpha - \lambda_{n\pm}] |n\rangle_e |g\rangle \}$ and the corresponding eigenvalues $\lambda_{n\pm} = \{ \Omega + (2n-1)\alpha \pm [(\Omega - \alpha)^2 + 4ng^2]^{1/2} \}/2$, where $\mathcal{N}_{n\pm}$ are the normalization constants. The schematic for the spectrum of JC model is shown in Fig. 1c. we notice that the eigenstates $|\lambda_{n\pm}\rangle$ are not orthogonal to each other. Thus, in the following calculations, we need to use the bi-orthogonal basis approach [17] with the eigenstates $|\lambda_{n\pm}^*\rangle$ of H_{eff}^* corresponding to the eigenvalues $\lambda_{n\pm}^*$. The orthogonal relations are $\langle \lambda_{n\mp}^* | \lambda_{n\pm} \rangle = 0$ and $\langle \lambda_{n\pm}^* | \lambda_{n\pm} \rangle = 1$.

In the above bi-orthogonal basis, we can evaluate the correlations $\langle \mathcal{T} a(t') a^\dagger(t) \rangle$ and G_4 , to obtain the single-photon and two-photon T -matrices and S -matrices. The results are listed as follows. **1.** For the bonding modes, the single photon S -matrix is $S_{pk} = t_k \delta_{kp}$, where $t_k = \exp(-i2\delta_k)$ and phase shift $\delta_k = \arg[(k - \lambda_{1+})(k - \lambda_{1-})]$. The anti-bonding modes are free of coupling, thus possess S -matrix element $S_{pk}^{(o)} = \delta_{kp}$. **2.** Then the reflection and transmission coefficients are obtained as $\bar{r}_k = (t_k - 1)/2$ and $\bar{t}_k = (t_k + 1)/2$, which agree with Ref. [16]. **3.** The two-photon T -matrix elements are explicitly obtained

$$iT_{p_1 p_2 k_1 k_2} = \frac{i\tilde{V}^4 g^4 (E - \alpha - \Omega) \delta_{p_1 + p_2, E}}{\pi \prod_{s=\pm} (E - \lambda_{2s})} \times \frac{[(E - 2\Omega)(E - 2\alpha) - 4g^2]}{\prod_{s=\pm} \prod_{i=1,2} (k_i - \lambda_{1s})(p_i - \lambda_{1s})}, \quad (7)$$

where $E = k_1 + k_2$ is the total energy of the incident photons. The square norm $T_2 = |T_{p_1 p_2 k_1 k_2}|^2$ of T -matrix element exhibits the two-photon background fluorescence in the bonding mode.

Two-photon wave functions — Let $a_R(x_1)$ ($a_L(x_2)$) denote the annihilation operators of right (left) moving

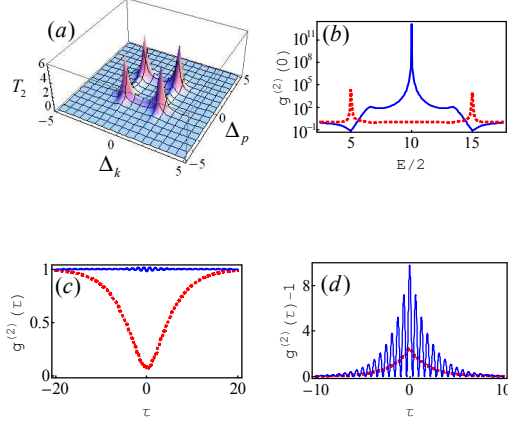


FIG. 2: (Color online) The two-photon background fluorescence and the correlation functions for strong coupling regime $g = 5$, the system parameters are $\omega_c = \Omega = 10$ and \tilde{V} is taken as unit: (a) Two-photon background fluorescence for $E = \lambda_{2s}^{(0)}$; (b) $g^{(2)}(0)$ of the reflected (solid (blue) curve) and transmitted (dashed (red) curve) photons; (c) $g^{(2)}(\tau)$ of the reflected (dashed (red) curve) photons for $E/2 = \lambda_{1s}^{(0)}$ and transmitted (solid (blue) curve) photons for $E/2 = \Omega$; (d) $(g^{(2)}(\tau) - 1)/10^{15}$ of the reflected (solid (blue) curve) photons for $E/2 = \Omega$ and $(g^{(2)}(\tau) - 1)/10^6$ of transmitted (dashed (red) curve) photons for $E/2 = \lambda_{1s}^{(0)}$.

photons [18]. It follows from Eqs. (3) and (7) that the out-going state $|X_{\text{out}}\rangle = |t_{\text{out}}\rangle + |r_{\text{out}}\rangle + |rt_{\text{out}}\rangle$ for two incident right-moving photons with momenta k_1 and k_2 contains three parts: (a) The quantum state of two transmitted photons

$$|t_{\text{out}}\rangle = \int dx_1 dx_2 t_2(x_1, x_2) a_L^\dagger(x_1) a_L^\dagger(x_2) |0\rangle |g\rangle$$

explicitly defined by the two-photon wavefunction

$$t_2(x_1, x_2) = \frac{1}{2\pi} e^{iEx_c} [\bar{t}_{k_1} \bar{t}_{k_2} \cos(\Delta_k x) - F(\lambda, x)], \quad (8)$$

where $\Delta_k = k_1 - k_2$,

$$F(\lambda, x) = \frac{\tilde{V}^4 g^4 \sum_{s=\pm} s(E - 2\lambda_{1s}) \exp[i(\frac{E}{2} - \lambda_{1,-s})|x|]}{4(\lambda_{1+} - \lambda_{1-}) \prod_{s=\pm} [(E - \lambda_{2s}) \prod_{i=1,2} (k_i - \lambda_{1s})]}, \quad (9)$$

and $x = x_1 - x_2$ and $x_c = (x_1 + x_2)/2$ are the relative and center of mass coordinates, respectively; (b) The quantum state of the two reflected photons

$$|r_{\text{out}}\rangle = \int dx_1 dx_2 r_2(x_1, x_2) a_L^\dagger(x_1) a_L^\dagger(x_2) |0\rangle |g\rangle$$

explicitly defined by

$$r_2(x_1, x_2) = \frac{1}{2\pi} e^{iEx_c} [\bar{r}_{k_1} \bar{r}_{k_2} \cos(\Delta_k x) - F(\lambda, x)]; \quad (10)$$

(c) the left-right entangled two-photon state

$$|rt_{\text{out}}\rangle = \int dx_1 dx_2 rt_2(x_1, x_2) a_L^\dagger(x_1) a_R^\dagger(x_2) |0\rangle |g\rangle$$

describes the scenario where one photon is transmitted while the other is reflected, where

$$rt_2 = \frac{1}{2\pi} e^{i\frac{E}{2}x} [(\bar{t}_{k_1} \bar{r}_{k_2} e^{2i\Delta_k x_c} + \bar{t}_{k_2} \bar{r}_{k_1} e^{-2i\Delta_k x_c}) - 2F(\lambda, 2x_c)]. \quad (11)$$

Photon statistics can be studied through the coherence functions [1] $g^{(2)}(\tau) = G^{(2)}(\tau)/|G^{(1)}(0)|^2$, where $G^{(1)}(\tau) = \langle F_{\text{out}} | a_F^\dagger(x + \tau) a_F(x) | F_{\text{out}} \rangle$, $G^{(2)}(\tau) = \langle F_{\text{out}} | a_F^\dagger(x) a_F^\dagger(x + \tau) a_F(x + \tau) a_F(x) | F_{\text{out}} \rangle$, and $F = R$ and L correspond to the transmitted photons and reflected photons respectively, while $|R_{\text{out}}\rangle = |t_{\text{out}}\rangle / \langle t_{\text{out}} | t_{\text{out}} \rangle$ and $|L_{\text{out}}\rangle = |r_{\text{out}}\rangle / \langle r_{\text{out}} | r_{\text{out}} \rangle$. We have $g^{(2)}(\tau) = C(\tau)/D$, where $C(\tau) = |t_2(x, x + \tau)|^2$ for photons in transmission, and $|r_2(x, x + \tau)|^2$ for photons in reflection, which is independent of x and D is the normalization constant.

Strong coupling regime. — To explore the physical consequence of the result above, we first consider the strong coupling regime with $g > \tilde{V}^2$. The two-photon background fluorescence T_2 and the correlation function $g^{(2)}(\tau)$ are shown in Fig. 2 on resonance, i.e., $\omega_c = \Omega$, and we assume the same energy for the two incident photons, i.e. $\Delta_k = 0$. When the average energy of two photon $E/2 = \lambda_{1\pm}^{(0)}$, T_2 has one sharp peak at $\Delta_k = \Delta_p = 0$, and the four peaks emerge (see Fig. 2a) for $E = \lambda_{2\pm}^{(0)}$.

In Fig. 2b, $g^{(2)}(0)$ as the functions of energy $E/2$ per photon are shown for two reflected and transmitted photons, respectively. The two reflected photons exhibit sub-Poissonian statistics ($g^{(2)}(0) \ll 1$) for $E/2 = \lambda_{1\pm}^{(0)}$, and super-Poissonian statistics ($g^{(2)}(0) \gg 1$) for $E/2 = \lambda_{2\pm}^{(0)}/2$ or Ω . For $E/2 = \lambda_{1\pm}^{(0)}$, the correlation functions $g^{(2)}(\tau)$ of two reflected (dashed (red) curve in Fig. 2c) and transmitted (dashed (red) curve in Fig. 2d) photons exhibit anti-bunching and bunching behaviors, respectively. When $E/2 = \Omega$, $g^{(2)}(\tau)$ of two reflected (solid (blue) curve in Fig. 2d) and transmitted (solid (blue) curve in Fig. 2c) photons exhibit large bunching and anti-bunching behaviors, respectively. Our results for the two reflected photons in the case of a side-coupled cavity exactly agree with that for the transmitted photons observed in the experiment [5] for the case of a directly coupled cavity as shown in Fig. 1b. This agreement is not accidental, since the transmitted photon states in the direct-coupled case can be mapped into the reflected photon states in the side-coupled case, using a transformation detailed in Ref. [16].

In the side-coupled structure considered here, the reflected light arises purely from the decaying amplitudes from the cavity. Thus, single photon reflection peaks at a single photon energy of $\lambda_{1\pm}^{(0)}$, which is the energy

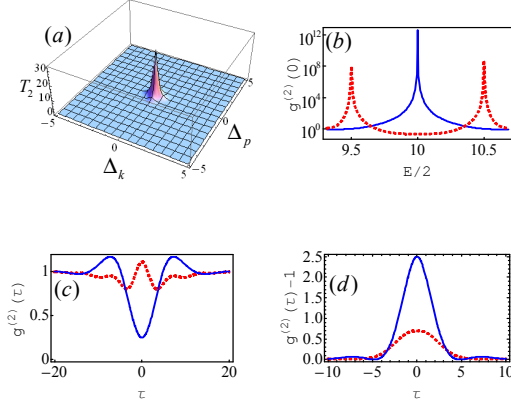


FIG. 3: (Color online) The two-photon background fluorescence and the correlation functions for weak coupling regime $g = 0.5$, the system parameters are the same as that in Fig. 2: (a)-(c) represent the same physical meaning in Fig. 2; (d) $(g^{(2)}(\tau) - 1)/10^{11}$ of the reflected (solid (blue) curve) photons for $E/2 = \Omega$ and $(g^{(2)}(\tau) - 1)/10^{10}$ of transmitted (dashed (red) curve) photons for $E/2 = \lambda_{1\pm}^{(0)}$.

level for one-photon dressed state in the cavity. Similarly, two-photon reflection peaks when the two photon energy is at $\lambda_{2\pm}^{(0)}$, where the cavity supports two-photon dressed states. However, since $\lambda_{2\pm}^{(0)} \neq 2\lambda_{1\pm}^{(0)}$, two photons each with energy $E/2 = \lambda_{1\pm}^{(0)}$ is off resonance from the two-photon dressed state. In such a case, the single excitation by the first photon in fact prevents the second photon from entering the cavity, resulting in the photon-blockade effect. Therefore, in contrast to the case with direct coupling where the photon-blockade effect manifests as a vanishing two-photon transmission, in the side-cavity case the photon blockade effect manifests as a vanishing two-photon reflection effect. When $E/2 = \Omega$, the single-photon reflection coefficient vanishes [16]. The two-photon reflection is due purely to the correlation induced by the TLS, which creates a two-photon bound state, and hence generates a large bunching effect [10]. We, therefore, for the first time, provide an exact analytic formula for the photon correlation function for this system, the special case of which agrees with the existing experimental data.

Weak coupling regime. — For the weak coupling regime with $g < \tilde{V}^2$, the two-photon background fluorescence T_2 and the correlation function $g^{(2)}(\tau)$ are shown in Fig. 3. Fig. 3a shows that the four peaks in Fig. 2a merge into a single peak when $E = \lambda_{2\pm}^{(0)}$.

We find that $g^{(2)}(0)$ of two reflected photons (Fig. 3b) has a simple structure with one peak at $E/2 = \Omega$ and always satisfies $g^{(2)}(0) \geq 1$, which means that the statistics of reflected photons can not be sub-Poissonian and the photon blockade effect vanishes due to the small energy splitting of $\lambda_{1\pm}^{(0)}$ for the weak coupling g . The bunching

behaviors exhibited by $g^{(2)}(\tau)$ of two reflected photons for $E/2 = \lambda_{1\pm}^{(0)}$ (dashed (red) curve in Fig. 3c) and Ω (solid (blue) curve in Fig. 3d) also illustrate the vanishing of photon blockade effect. In addition, $g^{(2)}(\tau)$ of anti-bunched and bunched transmitted photons for $E/2 = \Omega$ and $\lambda_{1\pm}^{(0)}$ are shown by the solid (blue) curve in Fig. 3c and the dashed (red) curve in Fig. 3d.

Conclusion. — We have analytically studied the two-photon transport in a waveguide coupled to the cavity containing TLS, and obtained an exact analytic formula describing the photon blockade effect. The quantum statistics of the outgoing photons are discussed in details using the exact two-photon wave-functions. These results agree with the observations of recent experiment [5]. Our theoretical approach can also be generalized to the many-photon transport cases.

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- [1] D. F. Walls and G. J. Milburn, *Quantum Optics* (Springer-Verlag, Berlin Heidelberg, 2008).
 - [2] H. J. Kimble, *Nature* **453**, 1023 (2008).
 - [3] A. Blais *et al.*, *Phys. Rev. A* **69**, 062320 (2004).
 - [4] F. Marquardt and S. M. Girvin, *Physics* **2**, 40 (2009); P. Zhang *et al.*, *Phys. Rev. Lett.* **95**, 097204 (2005); F. Xue *et al.*, *Phys. Rev. B* **75**, 033407 (2007).
 - [5] K. M. Birnbaum *et al.*, *Nature (London)* **436**, 87 (2005).
 - [6] T. Aoki *et al.*, *Nature* **443**, 671 (2006); T. Yoshie *et al.*, *Nature* **432**, 200 (2004); A. Badolato *et al.*, *Science* **308**, 1158 (2005); D. England *et al.*, *Nature* **450**, 857 (2007); K. Hennessy *et al.*, *Nature* **445**, 896 (2007); I. Fushman *et al.*, *Science* **320**, 769 (2008).
 - [7] A. Wallraff *et al.*, *Nature* **431**, 162 (2004); K. Srinivasan and O. Painter, *Nature* **450**, 862 (2007); B. Dayan *et al.*, *Science* **208**, 1062 (2008).
 - [8] J. T. Shen and S. Fan, *Phys. Rev. Lett.* **95**, 213001 (2005); *ibid.* **98**, 153003 (2007); *Opt. Lett.* **30**, 2001 (2005).
 - [9] L. Zhou *et al.*, *Phys. Rev. A* **76**, 012313 (2007); *Phys. Rev. Lett.* **101**, 100501 (2008); H. Dong *et al.*, *Phys. Rev. A* **76**, 063847 (2009).
 - [10] R. J. Brecha, P. R. Rice, and M. N. Xiao, *Phys. Rev. A* **59**, 2392 (1999).
 - [11] L. Tian and H. J. Carmichael, *Phys. Rev. A* **46**, 6801 (1992).
 - [12] E. Waks and J. Vuckovic, *Phys. Rev. Lett.* **96**, 153601 (2006); R. J. Thompson *et al.*, *Phys. Rev. Lett.* **68**, 1132 (1992).
 - [13] K. Srinivasan and O. Painter, *Phys. Rev. A* **75**, 023814 (2007).
 - [14] T. Shi and C. P. Sun, *Phys. Rev. B* **79**, 205111 (2009); arXiv:0907.2776.
 - [15] E. Rephaeli, J. T. Shen, and S. Fan, *Phys. Rev. A* **82**, 033804 (2010).
 - [16] J. T. Shen and S. Fan, *Phys. Rev. A* **79**, 023837 (2009).
 - [17] C. P. Sun, *Phys. Scr.* **48**, 393 (1993).
 - [18] J. T. Shen and S. Fan, *Phys. Rev. Lett.* **98**, 153003 (2007).

(2007); Phys. Rev. A **76**, 062709 (2007).